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On the Generation of Rogue Waves in Dusty Plasmas Due to Modulation Instability of Nonlinear Schrödinger Equation

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Abstract: The release of rogons (rogue waves) associated with the electrostatic perturbations in dusty plasmas containing two-temperature ions is investigated. Solving the fluid equations using perturbation method, a nonlinear Schrödinger equation (NLSE) is derived for the electrostatic potential amplitude, associated with the propagation of envelope wavepackets. The solution of the NLSE is presented, which proposed as an effective tool for studying the rogons in different Saturn's rings (i.e. E-ring, F-ring, and B-ring). Our analysis indicates that the plasma parameters of the E-ring cannot support the propagation of rogue waves, but the rogue waves may exist in B- and F-rings. The existence region of the created rogons is defined, as well as the forbidden zone of rogue waves is examined. The variation of the structural properties of the rogons with relevant plasma parameters is investigated, in particular focusing on the ratio between the low-temperature ion number density-to-dust number density, as well as the temperature ratio between the low-temperature ions-to-electrons.

Keywords: dusty plasma, rogue waves, Saturn's rings, nonlinear Schrödinger equation

1 Introduction

Dusty plasmas is a somewhat ambiguous term for the mixture of charged dust grains with the electrons and ions which found in normal plasmas. First, the Voyager observations during 1980s showed various phenomena in the Saturn's rings that could not be explained on purely gravitational alone, as well as a new ring (that is called Saturn's F-ring) discovered by these missions too. Furthermore, there are many examples in the solar system have relevance to the presence of dust grains such as in noctilucent clouds, interstellar dust clouds, cometary tails, planetary rings, solar nebula, etc. Dust particles are also found in the Earth environment such as production processes, flames, rocket exhausts, fusion devices, and many laboratory experiments (see e.g. Refs. [1,2]).

In general, dust grains are highly negatively (or positively) charged and massive grains in electron-ion plasma. So, the dust masses are responsible for the appearance of new types of waves and instabilities. One of these waves is the low frequency dust-acoustic wave (DAW), which was reported theoretically first by Rao et al. [3] and was verified experimentally by Barkan et al. [4]. Linear and nonlinear waves can exist in dusty plasmas based on the strength of the nonlinearity in the system. Many efforts have been made to examine different nonlinear modes depending on the tendency to self-organization and formation of long-living nonlinear dissipative and coherent structures in a dusty plasma, such as shock waves, solitons, cavitons, collapsing cavities, etc. (see e.g. Refs. [5,6,7,8,9,10,11,12,13,14, 15]). Both shocks and solitons in dusty plasmas can be formed by different means. These are not necessarily restricted to the mode excitation due to instabilities, or an external forcing, but can also be a regular collective process analogous to the shock wave generation in gas dynamics [2].

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One of the important effects that can change the nonlinear modes propagation is the temperature of the species, i.e. the presence of one species but having two-temperatures. For example, the ions could have low and high temperatures depending on achieving the condition of the two-temperatures ion assumption. The latter is defined by the energy rate E_R between the two type of ions must be much smaller than the characteristic frequency of the system $\omega_{pd} = (4\pi n_{d0} e^2 z_d^2/m_d)^{1/2}$, i.e., $E_R/\omega_{pd} \ll 1$, here ω_{pd} is the dust plasma frequency, n_{d0} is the unperturbed dust density, z_d is the dust charge number, and m_d is the dust mass. The two type ions are assumed to have the same mass m_i and charge number z_i, but one requires lower temperature T_{il} and another higher temperature T_{ih}. Here, $E_R = \Gamma/v_{ith}^2$, where $v_{ith} = \sqrt{T_{ih}/m_i}$and $\Gamma = (4\pi n_{il0} e^4 \ln\Lambda)/m_i^2$, with a Coulomb logarithm $\ln\Lambda \approx 10-15$ and unperturbed low temperature ion number density n_{il0} [16,17]. The validity of the two-temperature ions assumption is examined for Saturn's environments as in Refs. [18,19].

During the last two centuries, many reports of extreme wave events teem in ocean seafarer stories: an ultra-high ghost wave occurs unexpectedly, propagates for short times destroying everything in its passing and then disappears without a trace [20]. Now, these catastrophic waves are well-known by rogue waves (or freak waves, or monster waves, or rogons). On the other hand, rogons are short-lived phenomena appearing suddenly out of nowhere. The average height of the rogons can be two, three, or even more times the height of the surrounding waves.

Importantly, fundamental research has by now gone beyond the standard ocean-surface-dynamical problem, tracing rogue waves in different fields of science starting from mid-ocean and coastal waters [21,22], fiber optics [23], Bose-Einstein condensates [24,25], plasma physics (e.g. Refs. [26,27,28,29,30]), and even in finance [31]. Since the phenomenon of rogue waves is still a matter of active research, it is precocious to state clearly what the most common causes are or whether they vary from place to place. Several mechanisms for rogue waves including diffractive focusing, nonlinear effects (modulational instability), wind waves, and thermal expansion. Here, we are interesting to examine the generation of rogue waves based on one of the successful theories to explain the propagation of rogue waves which is the modulational instability due to the presence of nonlinear effects in the plasmas.

In view of the crucial importance of this challenging phenomenon, we have undertaken an investigation of the occurrence of rogue waves in different Saturn's rings which is associated with the presence of charged dust grains interacting with two-temperatures ions plasma. A reductive perturbation method is used to reduce the basis equations to one evolution equation describing the plasma system. The nonlinear Schrödinger equation is derived from the evolution equation, which admits a rational solution characterizing the rogue wave profile. Using the available data for different Saturn's rings, an ad hoc numerical study of the rogons is briefly presented.

2 Basic equations and formulation of the problem

Let us consider a system of four-components collisionless, unmagnetized dusty plasma consisting of negatively charged dust particles, isothermal electrons, and two-temperature isothermal ions. The dimensionless basic equations are governed by

$$\frac{\partial n_d}{\partial t} + \frac{\partial}{\partial x}(n_d u_d) = 0, \quad (1)$$

$$\left(\frac{\partial}{\partial t} + u_d \frac{\partial}{\partial x}\right) u_d - \frac{\partial \phi}{\partial x} = 0, \quad (2)$$

$$\frac{\partial^2 \phi}{\partial x^2} + n_{il} + n_{ih} - n_d - n_e = 0, \quad (3)$$

where the number densities of ions and electrons are expressed as:

$$n_{il} = \mu_{il}\ \exp(-\Delta_{il}\phi), \quad (4)$$

$$n_{ih} = \mu_{ih} \exp\left(-\Delta_{ih}\phi\right), \quad (5)$$

and

$$n_e = \mu_e\ \exp\left(\Delta_e \phi\right). \quad (6)$$

Here, u_d, n_d, and ϕ referred to the dust fluid velocity, number density, and electrostatic potential, respectively. The densities n_d and n_j ($j = il$, ih, and e) are normalized by n_{d0} and $n_d z_{d0}$, respectively. The space coordinate is normalized by the dust Debye length λ_{Dd} $=\left(T_{eff}/4\pi n_{d0} e^2 z_d\right)^{1/2}$, the time is normalized by the inverse of dust plasma frequency ω_{pd}^{-1} $=$ $\left(m_d/4\pi n_{d0} e^2 z_d^2\right)^{1/2}$, the velocity u_d is normalized by the dust-acoustic speed $C_{DA} = \left(z_d T_{eff}/m_d\right)^{1/2}$, and ϕ is normalized by $T_{eff}/e z_d$. At equilibrium, we have $\mu_{il} + \mu_{ih} = \mu_e + 1$, where

$$\frac{1}{T_{eff}} = \frac{1}{z_d n_{d0}}\left[\frac{n_{il0}}{T_{il}} + \frac{n_{ih0}}{T_{ih}} + \frac{n_{e0}}{T_e}\right],$$

$$\Delta_{il} = \frac{T_{eff}}{z_d T_{il}}, \qquad \Delta_{ih} = \frac{T_{eff}}{z_d T_{ih}}, \qquad \Delta_e = \frac{T_{eff}}{z_d T_e},$$

$$\mu_{il} = \frac{n_{il0}}{z_d n_{d0}}, \qquad \mu_{ih} = \frac{n_{ih0}}{z_d n_{d0}}, \qquad \mu_e = \frac{n_{e0}}{z_d n_{d0}}.$$

where T_e, T_{il}, and T_{ih} are the temperatures of electrons, cold ions, and hot ions in units of energy, respectively.

The electrons being more mobile than the other plasma particles and they will impact more into the grain surface, the grain will be negatively charged, acquire a negative potential with respect to the ambient plasma. After that, the random motion of the ions and the electrons in the neighborhood of the charged grains are disturbed. As the gains are negatively charged, the ions are attracted and electrons are repelled .Then, the grains become positively charged, the possibility of hitting ions decrease and the chance for electrons is increase. Finally, the electron flux is reduced by repulsion just enough to balance with ion flux. The dust grain charge q_d is determined by $dq_d/dt = \sum_j I_j$ where j represents the plasma species (electrons and ions) and I_j is the current associated with the species j. At balance the net current flowing onto the dust grain surface becomes zero, so we consider the dust charge is constant at this case.

To study small- but finite-amplitude DAWs, we derive the evolution equation describing the basic set of fluid equations (1)–(6). So, we employ the reductive perturbation method, which introduces the stretched space-time coordinates

$$X = \varepsilon^{1/2}(x - \lambda t) \quad \text{and} \quad T = \varepsilon^{3/2} t, \tag{7}$$

where λ is the phase speed of the acoustic waves that will be determined later and ε is a smallness parameter measures the weakness of the amplitude or dispersion. The physical quantities appear in Eqs. (1)–(6) are expanded as:

$$n_d = 1 + \varepsilon n_d^{(1)} + \varepsilon^2 n_d^{(2)} + \varepsilon^3 n_d^{(3)} + \ldots\ldots.., \tag{8}$$

$$u_d = \varepsilon u_d^{(1)} + \varepsilon^2 u_d^{(2)} + \varepsilon^3 u_d^{(3)} + \ldots\ldots\ldots.., \tag{9}$$

$$\phi = \varepsilon \phi^{(1)} + \varepsilon^2 \phi^{(2)} + \varepsilon^3 \phi^{(3)} + \ldots\ldots\ldots.., \tag{10}$$

Substituting Eqs. (7)–(10) into Eqs. (1)–(6), we obtain to the lowest order in ε

$$n_d^{(1)} = u_d^{(1)} = -\phi^{(1)}, \tag{11}$$

while the Poisson equation gives the compatibility condition

$$\lambda = (\mu_{il}\Delta_{il} + \mu_{ih}\Delta_{ih} + \mu_e\Delta_e)^{-1/2}. \tag{12}$$

To the next-order of ε, we get the system of equations in the second-order perturbed quantities. Solving those equations, we obtain the Korteweg-de Vries (KdV) equation

$$\frac{\partial \phi^{(1)}}{\partial T} + B\phi^{(1)}\frac{\partial \phi^{(1)}}{\partial X} + \frac{1}{2}\frac{\partial^3 \phi^{(1)}}{\partial X^3} = 0, \tag{13}$$

where

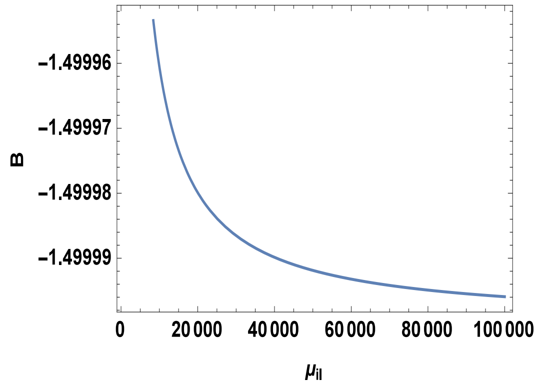

Fig. 1: The nonlinear coefficient B against μ_{il} for Saturn's E-ring, where $\beta_1 = 0.1$, $\beta_2 = 0.98$, $\beta = \beta_1/\beta_2$, and $\mu_{ih} = 1000$.

$$B = \frac{1}{2}\left[-3 + \frac{\mu_{il}(1-\beta_1^2) + \mu_{ih}\beta^2(1-\beta_2^2) + \beta_1^2}{(\mu_{il}(1+\beta_1) + \mu_{ih}\beta(1+\beta_2) - \beta_1)^2}\right],$$

$$\beta_1 = \frac{T_{il}}{T_e}, \qquad \beta_2 = \frac{T_{ih}}{T_e}, \qquad \text{and} \qquad \beta = \frac{T_{il}}{T_{ih}}.$$

It is well-known that the KdV equation (13) has different nonlinear solutions including solitary wave solution. However, the latter is out the scope of the present work since we are interesting to investigate the rogue wave solution of the evolution equation describing the plasma in different Saturn's rings. Describing the rogue waves could be done by using the nonlinear Schrödinger equation (NLSE). The KdV equation can be transformed to the NLSE for small wave number, however the coefficients of the produced NLSE cannot satisfy the condition for rogue wave existence. Therefore, the NLSE that has been obtained from the KdV equation cannot support rogue wave solution (see e.g. Ref. [32]). Now, it is interesting to examine the sign of the nonlinear coefficient B. On the other hand, could B equal zero at some critical value. Numerical analysis indicates that at critical concentration of low ions density (i.e. $\mu_{il} \equiv \mu_{ilc}$) the nonlinear coefficient $B = 0$. Thus, the KdV equation is not sufficient to describe the nonlinear wave propagation at μ_{ilc}, and we need to derive another equation describing the system. We can use new stretched variables as

$$X = \varepsilon(x - \lambda t) \quad \text{and} \quad T = \varepsilon^3 t, \tag{14}$$

along with the expansions (8)–(10) into the basic Eqs. (1)–(6), after some algebraic manipulations, we finally obtain the modified Korteweg-de Vries (mKdV) equation

$$\frac{\partial \phi^{(1)}}{\partial T} + C\phi^{(1)2}\frac{\partial \phi^{(1)}}{\partial X} + \frac{1}{2}\frac{\partial^3 \phi^{(1)}}{\partial X^3} = 0, \tag{15}$$

where

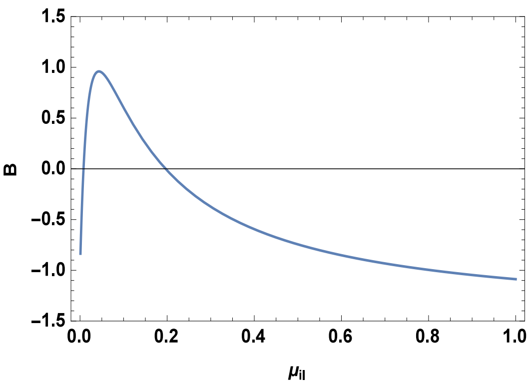

Fig. 2: The nonlinear coefficient B against μ_{il} for Saturn's B-ring, where $\beta_1 = 0.05$, $\beta_2 = 0.98$, $\beta = \beta_1/\beta_2$, and $\mu_{ih} = 1$.

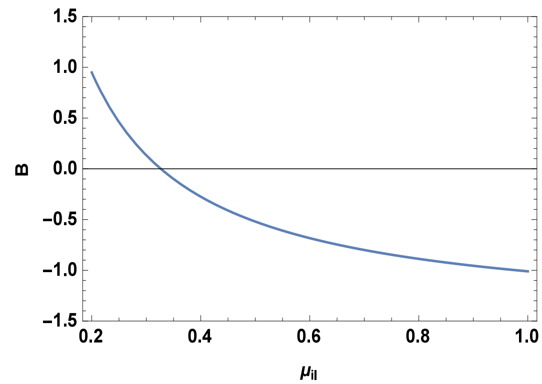

Fig. 3: The nonlinear coefficient B against μ_{il} for Saturn's F-ring, where $\beta_1 = 0.01$, $\beta_2 = 0.98$, $\beta = \beta_1/\beta_2$, and $\mu_{ih} = 0.8$.

$$C = \frac{1}{4}\left[15 - \frac{\mu_{ilc}(1+\beta_1^3) + \mu_{ih}\beta^3(1+\beta_2^3) - \beta_1^3}{(\mu_{ilc}(1+\beta_1) + \mu_{ih}\beta(1+\beta_2) - \beta_1)^3}\right].$$

Notice that we have replaced the expression of the phase velocity from Eq. (12) into the coefficients B and C, then make an appropriate simplification to ensure that the compatibility condition (12) is always satisfied.

It is interesting to transform the mKdV Eq. (15) to the NLSE to describe the behavior of the weakly nonlinear wave packet that gives rise to freak wave propagation. Therefore, we expand $\phi^{(1)}$ as [33, 34]

$$\varphi = \sum_{m=1}^{\infty} \varepsilon^m \sum_{L=-m}^{L=m} \varphi_L^{(m)}(X,T)\exp(iL\theta), \tag{16}$$

Here, $\varphi \equiv \varphi^{(1)}$ for simplicity and $\theta = (kX - \omega T)$, where k and ω are real variables representing the wave number and the frequency of the carrier wave, respectively. The stretched variables X and T are given by

$$\overline{X} = \varepsilon(X - v_g T) \quad \text{and} \quad \overline{T} = \varepsilon^2 T, \tag{17}$$

where v_g is the group velocity, $\varphi_L^{(m)}$ in equation (16) must be real, so we consider $\varphi_{-L}^{(m)} = \varphi_L^{(m)*}$, where the asterisk indicates the complex conjugate. Substituting Eqs. (16) and (17) into Eqs. (15) and comparing the coefficient of ε. We get in the first-order of the approximation for $m = 1$ and $L = 1$ the electrostatic waves dispersion relation $\omega = -\frac{1}{2}k^3$. The second order approximation for $m = 2$ and $L = 1$ yields $v_g = -\frac{3}{2}k^3$ is the group velocity. From the third-order approximation $m = 3$ and $L = 1$, we finally deduce from the resulting condition the NLSE equation which take the form:

$$i\frac{\partial\psi}{\partial\overline{T}} + \frac{1}{2}P\frac{\partial^2\psi}{\partial\overline{X}^2} + Q\,|\psi|^2\,\psi = 0. \tag{18}$$

Here, $\varphi_1^{(1)} \equiv \psi$ for simplicity that represents the electrostatic wave envelope. The coefficients of the dispersion and nonlinear terms are given by $P = -3k$ and $Q = -Ck$, respectively. The character of the dynamic wave depends on the sign of the ratio of $P/Q = 3/C$. The sign of this ratio refer to the (in)stability of the system. The unstable envelope pulses propagate when $(P/Q) > 0$, while the stable envelope pulses exist when $(P/Q) < 0$. The waves become stable if $C < 0$ and unstable if $C > 0$.

The NLSE (18) has a rational solution that is located on a nonzero background and localized both in the $\overline{T}$ and $\overline{X}$ directions as [35]

$$\psi = \sqrt{\frac{P}{Q}}\left[\frac{4(1+2iP\overline{T})}{1+4P^2\overline{T}^2+4\overline{X}^2} - 1\right]\exp(iP\overline{T}). \tag{19}$$

Equation (19) represents the freak wave solution in the unstable zone of the NLSE (18) for which the nonlinear coefficient C must be positive. If we substitute Eqs. (16) and (17) into the KdV equation (13), we obtain also the NLSE but with different forms of P and Q, i.e. $P = -3k$ and $Q = B^2/3k$. Therefore, the ratio P/Q is always negative and hence the NLSE that obtained from the KdV equation cannot support rogue wave solution and it is usualy represent stable wave.

3 Numerical analysis and discussion

Recalling that the transfer from KdV Eq. (13) to mKdV Eq. (15) needs a substantial condition which is the nonlinear coefficient B vanishes at critical concentration of low ions density (i.e. $\mu_{il} \equiv \mu_{ilc}$). To examine this condition for typical plasma parameters of E-, B-, and F-rings, we have plotted the nonlinear coefficient B against the concentration of low ions density μ_{il} as in Figs. 1, 2, and 3. It is seen that for E-ring typical parameters, the nonlinear coefficient B has always

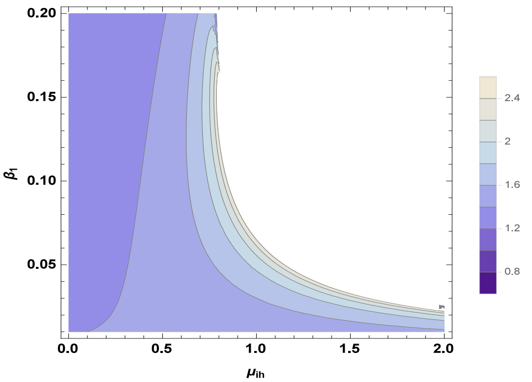

Fig. 4: The contour plot of B-ring of the rogue wave amplitude $\sqrt{P/Q}$ against μ_{ih} and β_1, where $\beta_2 = 0.98$ and $\beta = \beta_1/\beta_2$.

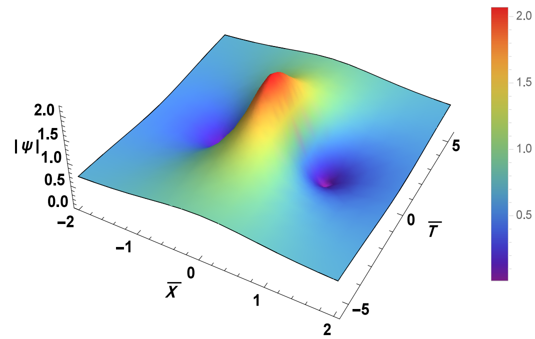

Fig. 6: The rogue pulse profile of B-ring, where $\mu_{ih} = 0.3$, $\beta_1 = 0.15$, $\beta_2 = 0.98$, and $\beta = \beta_1/\beta_2$.

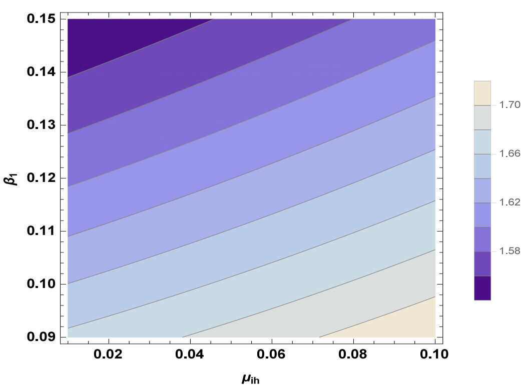

Fig. 5: The contour plot of F-ring of the rogue wave amplitude $\sqrt{P/Q}$ against μ_{ih} and β_1, where $\beta_2 = 0.98$ and $\beta = \beta_1/\beta_2$.

negative values. However, for B- and F-rings the nonlinear coefficient B has either positive or negative values, as well as at critical concentration of low ions density (i.e. $\mu_{il} \equiv \mu_{ilc}$) B vanishes, which is depicted in Figs. 2 and 3. Therefore, the transformation from the KdV to mKdV is only possible for B- and F-rings. Hence, the following analysis will consider only these two cases.

Recalling that the rogue waves can exist in the modulational instability region that represents by the NLSE (18). On the other hand, if $P/Q < 0$, the amplitude modulated envelope will be "stable" against external perturbations and the waves are modulationally "stable" and may propagate in the form of a "dark" ("black" or "gray") envelope wave packet, i.e., a propagating localized "hole" (a "void") amidst a uniform wave energy region. In other words, for "positive" P/Q, the carrier wave is modulationally "unstable;" it may either "collapse," due to (possibly random) external perturbations, or lead to the formation of "bright" envelope modulated wave packets, i.e., localized envelope "pulses" confining the carrier wave. We are interesting to investigate a special modulational unstable solution for $P/Q > 0$, which is local both in space and time. This is typical for rogue waves, which appear suddenly and then disappear without trace. In order to gain some insight, we have depicted the ratio P/Q for B- and F-rings. From Eq. (19), it is seen that the maximum amplitude of the rogons is proportional to $(P/Q)^{1/2} = (3/C)^{1/2}$. So, for rogons existence the sign of the nonlinear coefficient C should be positive. The contour plot of the ratio $(3/C)^{1/2}$ is depicted in Figs. 4 and 5 for B- and F-rings, respectively. Figure 4 clears the contour plot of the ratio $(3/C)^{1/2}$ for B-ring. It is obvious that the rogons can exist for wide range of plasma parameters except for the white zone since the ratio P/Q is negative or $C < 0$. The enhancement of the parameter β_1 would lead to shrink (increase) the maximum amplitude for $\mu_{ih} < 0.5$ ($\mu_{ih} > 0.5$), however the parameter μ_{ih} increases the pulse amplitude. The behavior in the F-ring is different, since the rogue waves can exist for the possible values of plasma parameters as depicted in Fig. 5. Now, we will plot the rogue wave profile for different Saturn's rings as shown in Figs. 6 and 7. It is seen that the rogue wave amplitude for B-ring is shorter than F-ring. Furthermore, the rogue wave profiles have not change in their spatial size. The taller pulse amplitude in the F-ring may be speculated to that the plasma parameters in the F-ring produce sufficient high nonlinearity to absorb a large amount of energy from the background waves and produce towering rogons.

4 Summary

We have presented a theoretical model, for rogons (rogue waves) associated with the electrostatic pulse propagation in dusty plasmas containing two-temperature ions. Solving the fluid equations, we have derived the nonlinear

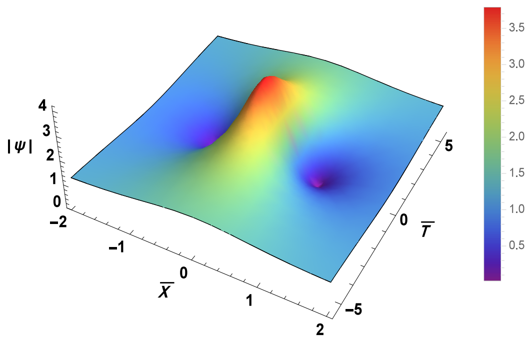

Fig. 7: The rogue pulse profile for F-ring, where $\mu_{ih} = 0.05$, $\beta_1 = 0.13$, $\beta_2 = 0.98$, and $\beta = \beta_1/\beta_2$.

Schrödinger equation (NLSE) for the electrostatic potential amplitude, associated with the propagation of envelope wave packets. An envelope rational solution of the NLSE is presented, which proposed as an effective tool for studying the rogons in Saturn's rings. Three rings namely; E-ring, F-ring, and B-ring are considered, and the existence of rogue waves in these rings are examined. It is found that only B- and F-rings support the propagation of rogue waves. The existence region of the rogons is examined for typical plasma parameters of B- and F-rings, and the forbidden region of the rogons are defined precisely. Outside the forbidden region, it is possible for a random perturbation to grow and may thus lead to the creation of the rogue waves. We can elucidate this phenomenon to that the relatively high number density of the dust grains in the B- and F-rings could create an abundant nonlinearity in the system to soak up high amount of energy from the background waves producing steep rogons.

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